DOE/ER/40762-274
UM-PP#03-035

Counting Rule for Hadronic Light-Cone Wave Functions

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(Dated: June 4, 2005)

Abstract

We introduce a systematic way to write down the Fock components of a hadronic light-cone wave function with n partons and orbital angular momentum projection l_z. We show that the wave function amplitude $\psi_n(x_i, k_{i\perp}, l_{zi})$ has a leading behavior $1/(k_\perp^2)^{[n+|l_z|+\min(n'+|l'_z|)]/2-1}$ when all parton transverse momenta are uniformly large, where n' and l'_z are the number of partons and orbital angular momentum projection, respectively, of an amplitude that mixes under renormalization. The result can be used as a constraint in modeling the hadronic light-cone wave functions. We also derive a generalized counting rule for hard exclusive processes involving parton orbital angular momentum and hadron helicity flip.

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Light-cone wave functions are useful tools to describe physics of hadrons in high-energy scattering. They are snap-shots of hadrons when the latter are moving with the speed of light (infinite momentum frame). These wave functions can be obtained, in principle, through solving the eigen-equation of the light-cone Hamiltonian using either analytical or numerical methods [1, 2]. They can also be obtained from the Bethe-Salpeter amplitudes by integrating out the k^- components of the parton four-momenta if the latter are known. Their moments in momentum space can be calculated using lattice QCD or the QCD sum-rule methods [3, 4]. In phenomenological approaches, the light-cone wave functions are parametrized to fit experiment data [5, 6, 7, 8].

In Ref. [9], we have proposed a systematic way to enumerate independent amplitudes of a light-cone wave function by writing down the matrix elements of a class of light-cone-correlated quark-gluon operators, in much the same way that has been used to construct independent light-cone distribution amplitudes in which the parton transverse momenta are integrated over [10]. In Ref. [11], we have applied this approach to the nucleon, finding that six amplitudes are needed to describe the three-quark sector of the nucleon wave function.

In this paper, we introduce a direct method of constructing the light-cone wave functions in momentum space. By exploiting the relations between light-cone amplitudes and the matrix elements of light-cone-correlated quark-gluon operators, we study how the wave function amplitudes depend on the transverse momenta of partons in the asymptotic limit. We find that a general amplitude $\psi_n(x_i, k_{\perp i}, \lambda_i, l_{zi})$ describing an n-parton state with orbital angular momentum projection l_z goes like

$$\psi_n(x_i, k_{\perp i}, \lambda_i, l_{zi}) \rightarrow 1/(k_\perp^2)^{[n+|l_z|+\min(n'+|l_z'|)]/2-1} , \tag{1}$$

in the limit that $k_{1\perp} \sim k_{2\perp} \sim ... \sim k_{n-1\perp} \sim k_\perp \rightarrow \infty$, where n' and l_z' characterize the amplitude that mixes under scale evolution. The result explains the scaling behavior of the $F_2(Q^2)$ form factor obtained recently in perturbative QCD [12], and helps to establish more general scaling properties of exclusive scattering amplitudes [13, 14, 15, 16, 17]. It also can be used as a constraint in building phenomenological wave functions of hadrons consistent with perturbative QCD.

Let us first introduce a systematic method to construct the light-cone Fock wave function of a hadron with helicity Λ. Suppose a Fock component has n partons with creation operators $a_1^\dagger$, ..., $a_n^\dagger$, where the partons can either be gluons or quarks and the subscripts label the partons' quantum numbers such as spin, flavor, color, momentum, etc. Assume all color, flavor (for quarks) indices have been coupled properly using Clebsch-Gordon coefficients. The longitudinal momentum fractions of the partons are x_i $(i = 1, 2, ..., n)$, satisfying $\sum_{i=1}^n x_i = 1$, and the transverse momenta $k_{1\perp}, ..., k_{n\perp}$, satisfying $\sum_i^n \vec{k}_{i\perp} = 0$. We will eliminate $\vec{k}_{n\perp}$ in favor of the first $n-1$ transverse momenta. Assume the orbital angular momentum projections of the partons are $l_{z1}, ..., l_{z(n-1)}$, respectively, and let $l_z = \sum_{i=1}^{n-1} l_{zi}$, then

$$l_z + \lambda = \Lambda , \tag{2}$$

where $\lambda = \sum_{i=1}^n \lambda_i$ is the total parton helicity. Without loss of generality, we assume $l_z \geq 0$; even then, l_{zi} can have both signs. Thus, a general term in the hadron wave function appears as

$$\int \prod_{i=1}^n d[i] \;\; (k_{1\perp}^\pm)^{|l_{z1}|} (k_{2\perp}^\pm)^{|l_{z2}|} ... (k_{(n-1)\perp}^\pm)^{|l_{z(n-1)}|} \; \psi_n(x_i, k_{\perp i}, \lambda_i, l_{zi}) \; a_1^\dagger a_2^\dagger ... a_n^\dagger |0\rangle , \tag{3}$$

where $k_i^{\pm} = k_{ix} \pm k_{iy}$ and the $+(-)$ sign applys when l_{zi} is positive (negative), and $d[i] = dx_i d^2k_{\perp i}/(\sqrt{2x_i}(2\pi)^3)$ with the overall constraint on x_i and $k_{\perp i}$ implicit.

The above form can be further simplified as follows. Assume l_{zi} is positive and l_{zj} negative, and $l_{zi} > |l_{zj}|$, we have

$$\begin{aligned}
(k_i^+)^{l_{zi}}(k_j^-)^{-l_{zj}} &= (k_i^+)^{l_{zi}+l_{zj}}(k_i^+k_j^-)^{-l_{zj}} \\
&= (k_i^+)^{l_{zi}+l_{zj}}(k_{i\perp}\cdot k_{j\perp} + i\epsilon^{\alpha\beta}k_{i\alpha}k_{j\beta})^{-l_{zj}} \\
&= (k_i^+)^{l_{zi}+l_{zj}}\left(\phi_0 + \phi_1 i\epsilon^{\alpha\beta}k_{i\alpha}k_{j\beta}\right) \ , \qquad (4)
\end{aligned}$$

where $\alpha, \beta = 1, 2$, $\phi_{0,1}$ are polynomials in $k_{i\perp}^2$, $k_{j\perp}^2$, and $k_{i\perp}\cdot k_{j\perp}$. On the last line of the above equation we have used the identity $\epsilon^{\alpha\beta}\epsilon^{\gamma\delta} = \delta^{\alpha\gamma}\delta^{\beta\delta} - \delta^{\alpha\delta}\delta^{\beta\gamma}$. If $l_{zi} + l_{zj} \neq 0$, one can use $i\epsilon^{\alpha\beta}k_{1\alpha}k_{2\beta}k_1^+ = k_{1\perp}\cdot k_{1\perp}k_2^+ - k_{1\perp}\cdot k_{2\perp}k_1^+$ to further reduce the second term in the bracket. Following the above procedure, we can eliminate all negative l_{zj}, a general $l_z > 0$ component in the wave function reads

$$\begin{aligned}
&\int \prod_{i=1}^n d[i] \quad (k_{1\perp}^+)^{l_{z1}}(k_{2\perp}^+)^{l_{z2}}...(k_{(n-1)\perp}^+)^{l_{z(n-1)}} \\
&\quad \times \left(\psi_n(x_i, k_i, \lambda_i, l_{zi}) + \sum_{i<j=1|_{l_{zi}=l_{zj}=0}}^{n-1} i\epsilon^{\alpha\beta}k_{i\alpha}k_{j\beta}\psi_{n(ij)}(x_i, k_{\perp i}, \lambda_i, l_{zi})\right) \ a_1^\dagger a_2^\dagger ... a_n^\dagger|0\rangle \qquad (5)
\end{aligned}$$

where $\sum_i l_{zi} = l_z$ and $l_{zi} \geq 0$, and the sums over i and j are restricted to the $l_{zi} = 0$ partons.

Using the above construction, it is easy to see that the proton state with three valence quarks has six independent scalar amplitudes $\psi_{uud}^{(i)}, i = 1, ..., 6$ [11]. The wave function amplitudes for three quarks plus one gluon will be presented in a separate publication [18].

The mass dimension of ψ_n can be determined as follows: Assume the nucleon state is normalized relativistically $\langle P|P'\rangle = 2E(2\pi)^3\delta^3(\vec{P}' - \vec{P})$, $|P\rangle$ has mass dimension -1. Likewise, the parton creation operator $a_i^\dagger$ has mass dimension -1. Given these, the mass dimension of ψ_n is $-(n + |l_z| - 1)$. The mass dimension of $\psi_{n(ij)}$, however, is $-(n + |l_z| + 1)$ which can be accounted for by the previous formula with an effective angular momentum projection $|l_z| + 2$.

To find the asymptotic behavior of an amplitude $\psi_n(x_i, k_i, l_{zi})$ in the limit that all transverse momenta are uniformly large, we consider the matrix element of a corresponding quark-gluon operator between the QCD vaccum and the hadron state

$$\langle 0|\phi_{\mu_1}(\xi_1)....\phi_{\mu_n}(\xi_n)|P\Lambda\rangle \ , \qquad (6)$$

where ϕ are parton fields such as the "good" $(+)$ components of quark fields or $F^{+\alpha}$ of gluon fields, and μ_i are Dirac and transverse coordinate indices when appropriate. All spacetime coordinates ξ_i are at equal light-cone time, $\xi_i^+ = 0$. Fourier-transforming with respect to all the spatial coordinates $(\xi_i^-, \xi_{i\perp})$, we find the matrix element in the momentum space, $\langle 0|\phi_{\mu_1}(k_1)....\phi_{\mu_{n-1}}(k_{n-1})\phi_{\mu_n}(0)|p\Lambda\rangle \equiv \psi_{\mu_1,...,\mu_n}(k_1, ..., k_{n-1})$, here we have just shown $n-1$ parton momenta because of the overall momentum conservation. The matrix element can be written as a sum of terms involving projection operator $\Gamma^A_{\mu_1...\mu_n}(k_{\perp i})$ multiplied by scalar amplitude $\psi_{nA}(x_i, k_{\perp i}, l_{zi})$:

$$\begin{aligned}
\langle 0|\phi_{\mu_1}(k_1)....\phi_{\mu_{n-1}}(k_{n-1})\phi_{\mu_n}(0)|p\Lambda\rangle &\equiv \psi_{\mu_1,...,\mu_n}(k_1, ..., k_{n-1}) \\
&= \sum_A \Gamma^A_{\mu_1...\mu_n}(k_{\perp i})\psi_n^{(A)}(x_i, k_{\perp i}, l_{zi}) \ , \qquad (7)
\end{aligned}$$

where the projection operator Γ^A contains Dirac matrices and is a polynomial of order l_z in parton momenta. For example, the two quark matrix element of the pion can be written as [9],

$$\begin{aligned}&\langle 0|\overline{d}_{+\mu}(0)u_{+\nu}(x,k_\perp)|\pi^+(P)\rangle \\ &= (\gamma_5 \not{P})_{\nu\mu}\psi^{(1)}_{u\overline{d}}(x,k_\perp,l_z=0) + (\gamma_5\sigma^{-\alpha})_{\nu\mu}P^+k_{\perp\alpha}\psi^{(2)}_{u\overline{d}}(x,k_\perp,l_z=1)\ ,\end{aligned} \tag{8}$$

where the projection operators are shown manifestly. More examples for the proton matrix elements can be found in Ref. [11].

The matrix element of our interest is, in fact, a Bethe-Salpeter amplitude projected onto the light cone. One can write down formally a Bethe-Salpeter equation which includes mixing contributions from other light-cone matrix elements. In the limit of large transverse momentum $k_{\perp i}$, the Bethe-Salpeter kernels can be calculated in perturbative QCD because of asymptotic freedom. In the lowest order, the kernels consist of a minimal number of gluon and quark exchanges linking the active partons. For the lowest Fock components of the pion wave function, one gluon exchange is needed to get a large transverse momentum for both quarks [16]. As we shall see, asymptotic behavior of the wave function amplitudes depends on just the mass dimension of the kernels.

Schematically, we have the following equation for the light-cone amplitudes,

$$\begin{aligned}\psi_{\alpha_1,...,\alpha_n}(k_1,...,k_{n-1}) &= \sum_A \Gamma^A_{\alpha_1...\alpha_n}(k_{\perp i})\psi^A_n(x_i,k_{\perp i},l_{zi}) \\ = \sum_{n',\beta_1,...,\beta_{n'}}&\int d^4q_1...d^4q_{n'-1}H_{\alpha_1...,\alpha_n,\beta_1,...,\beta_{n'}}(q_i,k_i)\psi_{\beta_1,...,\beta_{n'}}(q_1,...,q_{n'-1})\end{aligned} \tag{9}$$

where $H_{\alpha_1,...,\alpha_n,\beta_1,...,\beta_{n'}}$ are the Bethe-Salpeter kernels multiplied by the parton propagators. When the parton transverse momenta are uniformly large, the kernels can be approximated by a sum of perturbative diagrams. The leading contribution to the amplitudes on the left can be obtained by iterating the above equation, assuming the amplitudes under the integration sign contain no hard components. As such, the integrations over $q_{\perp i}$ can be cut-off at a scale μ where $k_\perp \gg \mu \gg \Lambda_{\rm QCD}$, and the q_i dependence in H can be expanded in Taylor series. In order to produce a contribution to $\psi^{(A)}_n(x_i,k_{\perp i},l_{zi})$, the hard kernels must contain the projection operator $\Gamma^A_{\alpha_1...\alpha_n}(k_1,...,k_{n-1})$. Hence we write

$$\begin{aligned}&H_{\alpha_1...,\alpha_n,\beta_1,...,\beta'_n}(q_i,k_i) \\ &= \sum_{A,B}\Gamma^A_{\alpha_1...\alpha_n}(k_{\perp i})H_{AB}(x_i,k_{\perp i},y_i)\Gamma^B_{\beta_1...\beta_{n'}}(q_{\perp i})\ ,\end{aligned} \tag{10}$$

where $\Gamma^B_{\beta_1...\beta_{n'}}(q_{\perp i})$ is again a projection operator and $H_{AB}(x_i,k_i,y_i)$ are scalar functions of the transverse momenta $k_{\perp i}$ invariants. Substituting the above into Eq.(9) and integrating over q_i^-,

$$\begin{aligned}&\psi^{(A)}_n(x_i,k_{\perp i},l_{zi}) \\ &= \sum_{B,\beta_i}\int dy_1...dy_{n'-1}H_{AB}(x_i,k_i,y_i)\int d^2q_{\perp 1}...d^2q_{\perp(n'-1)}\Gamma^B_{\beta_1...\beta_{n'}}(q_{\perp i})\psi_{\beta_1,...,\beta_{n'}}(y_i,q_i) \\ &= \sum_{B,\beta_i,A'}\int dy_1...dy_{n'-1}H_{AB}(x_i,k_i,y_i)\int d^2q_{\perp 1}...d^2q_{\perp(n'-1)}\Gamma^B_{\beta_1...\beta_{n'}}(q_{\perp i}) \\ &\quad\times\Gamma^{A'}_{\beta_1...\beta_{n'}}(q_{\perp i})\psi^{(A')}_{n'}(y_i,q_{\perp i},l'_{zi})\ ,\end{aligned} \tag{11}$$

where the integrations over $q_{\perp i}$ are non-zero only when the angular momentum content of Γ^B and $\Gamma^{A'}$ is the same. Now the large momenta $k_{\perp i}$ are entirely isolated in H_{AB} which does not depend on any soft scale. The asymptotic behavior of $\psi_n^{(A)}(k_{\perp i})$ is determined by the mass dimension of H_{AB}, which can be obtained, in principle, by working through one of the simplest perturbative diagrams.

A much simpler way to proceed is to use light-cone power counting in which the longitudinal mass dimension, such as P^+, can be ignored because of the boost invariance of the above equation along the z direction. We just need to focus on the transverse dimensions. Since the mass dimension of the amplitudes is $-(n+|l_z|-1)$, that of $\Gamma^B\Gamma^{A'}$ is $2|l'_z|$, and the integration measure $2(n'-1)$, a balance of the mass dimension yields $[H_{AB}] = -(n-1+|l_z|)-(n'-1+|l'_z|)$. Therefore, we arrive at the central result of the paper that the leading behavior of the wave function amplitude goes as

$$\psi_n^{(A)}(x_i, k_{\perp i}, l_{zi}) \sim \frac{1}{(k_\perp^2)^{[n+|l_z|+\min(n'+|l'_z|)]/2-1}} \ , \tag{12}$$

which is determined by a mixing amplitude with smallest $n'+|l'_z|$. Since the wave function has mass dimension of $-(n+|l_z|-1)$, the coefficient of the asymptotic form must have a soft mass dimension $\Lambda_{\rm QCD}^{\min(n'+|l'_z|)-1}$.

For the quark-antiquark amplitudes of the pion, the leading behavior is determined by self-mixing: $\psi^{(1)}_{u\bar d}(x, k_\perp, l_z=0) \sim 1/k_\perp^2$ and $\psi^{(2)}_{u\bar d}(x, k_\perp, l_z=1) \sim 1/(k_\perp^2)^2$. On the other hand, for the three-quark amplitudes of the proton [11], we have, $\psi^{(1)}_{uud}(x_i, k_{\perp i}) \sim 1/(k_\perp^2)^2$, $\psi^{(2,3,4,5)}_{uud}(x_i, k_{\perp i}) \sim 1/(k_\perp^2)^3$, $\psi^{(6)}_{uud}(x_i, k_{\perp i}) \sim 1/(k_\perp^2)^4$. Here we recall that for $\psi^{(2)}_{uud}$, the effective angular momentum projection is $l_z^{\rm eff}=2$. Its leading behavior is determined by its mixing with $\psi^{(1)}_{uud}$.

What are the selection rules for amplitude mixings? First of all, because of angular momentum conservation, wave function amplitudes belonging to different hadron helicity states do not mix. Second, because of the vector coupling in QCD, the quark helicity in a hard process does not change. Therefore, the pion amplitude $\psi^{(2)}_{u\bar d}$ does not mix with $\psi^{(1)}_{u\bar d}$ because the total quark helicity differs. An example of the nontrivial amplitude mixing is between the pion's two-quark-one-gluon and two-quark amplitudes. If one calculates the asymptotic pion form factor using the wave function amplitudes directly, the three-parton component does contribute at the leading order. If, however, the form factor is calcualted using a factorization approach in which the amplitudes are only used at a soft-scale $\Lambda_{\rm QCD}$, the three parton compoent contributes as a higher twist.

As an example, we apply the amplitude counting rule to hard exclusive processes in which the leading light-cone wave functions of participating hadrons dominate. One can, of course, use the light-cone wave functions to calculate directly hard scattering amplitudes and cross sections, finding the asymptotic behavior of these physics observables. Using the expression derived for $F_2(Q^2)$ in Ref. [11] and the asymptotic behavior of $\psi^{(1)}_{uud} \sim 1/k_\perp^4$ and $\psi^{(3,4)}_{uud} \sim 1/k_\perp^6$, we easily derive the result found in Ref. [12]:

$$F_2(Q^2) \sim 1/(Q^2)^3 \sim F_1(Q^2)/Q^2 \ , \tag{13}$$

in asymptotic limit. On the other hand, the proton amplitudes $\psi^{(3,4)}_{uud}$ obtained from Melosh rotation are suppressed by only one power of $k_\perp$ relative to $\psi^{(1)}_{uud}$, and are inconsistent with perturbative QCD in the large $k_\perp$ limit [19]. It seems, however, that the Melosh-rotated

$\psi_{uud}^{(3,4)}$ amplitudes with a harder $k_\perp$-dependence are phenomenologically interesting to model $F_2(Q^2)/F_1(Q^2) \sim 1/Q$ behavior at an intermediate Q^2 [7, 20, 21, 22].

A simpler way to find a generalized counting rule for hard exclusive processes [13, 14] is to count the soft mass dimensions in scattering amplitudes; the scaling in hard kinematic variables is then determined by dimensional balance. For example, the wave function amplitude $\psi_n(x_i, k_i, l_{zi})$ when used in a factorization formula contains a soft scale factor $\Lambda_{QCD}^{n+|l_z|-1}$. Therefore a scattering amplitude involving $H = 1, ..., N$ hadrons with the light-cone amplitudes $\psi_{n_H}(x_i, k_i, l_{zi})$ contains a soft mass factor $\Lambda_{QCD}^{\sum_H(n_H+|l_{zH}|-1)}$. In the hadronic process $A+B \to C+D+...$, the fixed-angle scattering cross section calculated using the amplitudes $\psi_n(x_i, k_i, l_{zi})$ goes like

$$\Delta\sigma \sim s^{-1-\sum_H(n_H+|l_{zH}|-1)} , \tag{14}$$

where H sums over all hadrons involved. For $l_{zH} = 0$ and minimal n, this is just the counting rule of Brodsky-Farrar [13] and Matveev-Muradian-Tavkhelidze [14]. The derivation here emphasizes that the traditional counting rule applies only to hadron helicity conserving processes [17]. The generalized counting rule here applies to any hard process proceeded through any wave function amplitudes. In particular, it reproduces the result of Chernyak and Zhitnitsky for form factors where parton orbital angular momentum was first considered [15].

As an application, we consider pp elastic scattering. Three helicity conservation amplitudes are known to go like $M(++ \to ++) \sim M(+- \to +-) \sim M(-+ \to +-) \sim 1/s^4$ [17]. Our counting rule provides the scaling behavior of the helicity flipping amplitudes $M(++ \to +-) \sim 1/s^{9/2}$ and $M(-- \to ++) \sim 1/s^5$.

We end the paper with a few cautionary notes. First, we have ignored the Lanshoff type of contributions in hadron-hadron scattering [23]. Second, in an actual calculation of a scattering amplitude, there are integrations over partons' light-cone fractions x_i. These integrations may be divergent at the endpoints $x_i = 0, 1$ depending upon the choices of the light-cone wave functions. The QCD factorization and the naive power counting break down there [24, 25]. Finally, the light-cone wave functions defined in the light-cone gauge have singularities [26]. When regularized, Sudakov type of form factors appear which lead to the dependence of the light-cone wave functions on P^+ [27]. The $k_\perp$ counting breaks down in the region where the Sudakov form factors are important. However, in certain cases the endpoint singularities are regulated by the Sudakov effects, and the last two adverse factors cancel [28], leaving the naive counting rule intact. It is not clear, however, that this happens in general.

X. J. and F. Y. were supported by the U. S. Department of Energy via grant DE-FG02-93ER-40762. J.P.M. was supported by National Natural Science Foundation of P.R. China through grand No.19925520.

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